

Unquenching the Quark Model and Screened Potentials

Eric S. Swanson^{*1}

¹ Rudolph Peierls Centre for Theoretical Physics, Oxford University, Oxford, UK, OX1 3NP.

The low-lying spectrum of the quark model is shown to be robust under the effects of ‘unquenching’. In contrast, the use of screened potentials is shown to be of limited use in models of hadrons. Applications to unquenching the lattice Wilson loop potential and to glueball mixing in the adiabatic hybrid spectrum are also presented.

I. INTRODUCTION

Lattice gauge Wilson loop computations convincingly demonstrate that the interquark interaction grows linearly with distance. However, these computations are, somewhat paradoxically, made in the pure gauge theory: colour sources and sinks are taken to be static and virtual quarks are omitted. It is expected that the introduction of light quarks will allow the chromoelectric flux tube to break thereby screening the Wilson loop confinement potential at a range given by $br = 2M_{qQ} - 2M$ where b is the lattice string tension, M_{qQ} is the mass of a light quark – static source system, and M is the mass of the heavy quark.

The constituent quark model is similar to quenched lattice gauge theory in the sense that nonvalence quark effects are neglected (or are absorbed into parameters) in both formalisms. Incorporating quark loop (for example, meson creation) effects in the quark model has been a longstanding goal of hadronic physics[1]. It is therefore natural to enquire whether the use of screened potentials in the constituent quark model can account for some of these effects. Indeed several papers have appeared recently which use this idea and which claim good agreement with experiment[2].

More fundamentally, does the constituent quark model survive unquenching? It is possible that mass shifts induced by meson loops will severely distort the spectrum. I shall demonstrate that this need not be the case with the aid of a simple example below. It is also argued that the use of screened potentials is not a well-founded approach to unquenching the quark model. An obvious problem is that the bound state spectrum disappears above the ‘ionization’ energy, $2M_{qQ}$. While one may regard this as satisfactory if hadrons become so broad as to be unrecognizable above the ionization energy, there is no evidence of this in the spectrum. Furthermore, the prospect of free quarks emerging at higher energies is not in keeping with the current understanding of confinement.

A defining characteristic of the quark potential model is that the number of constituents in any hadronic sys-

tem is fixed. How this property is modified in a field theoretic context can be illustrated by considering the interaction between electrons in lattice QED. In this case, the quenched Wilson loop interaction may be interpreted as a potential because the imposition of the infinite electron mass limit removes all non-instantaneous interactions from the theory. The resulting interaction is then simply $V_{WL} = \alpha/r$. Unquenching QED introduces light fermion loop corrections to the instantaneous interaction (and light-by-light scattering effects) and this interaction can no longer be interpreted as a potential. Specifically, as the momentum transfer increases, the virtual fermion pair can go on shell and the number of constituents has changed. Increasing the momentum transfer further creates highly virtual e^+e^- pairs, and the potential interpretation of the interaction is regained.

The generic situation is shown in Fig. 1. It is typical to focus attention on the lowest adiabatic surface, which has suffered obvious modification. However, higher surfaces remain and, as the dashed line in the figure indicates, the concept of a $q\bar{q}$ potential is still possibly useful. This will be demonstrated in the next section.

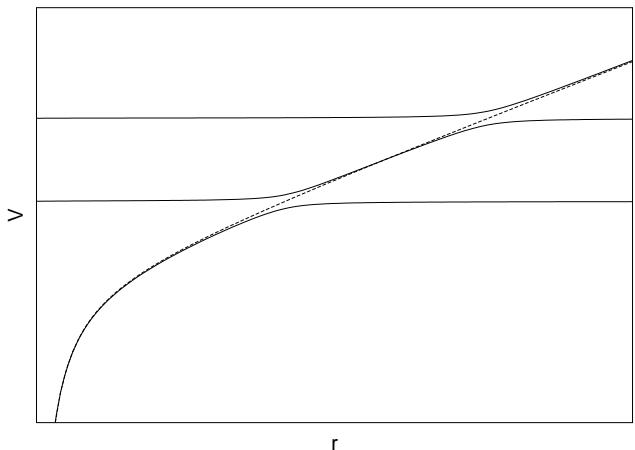


FIG. 1: Adiabatic Surface Mixing.

*on leave from the Department of Physics and Astronomy, University of Pittsburgh, Pittsburgh PA 15260.

II. A SIMPLE MODEL OF UNQUENCHING

A. Model Definition

It is convenient to define a simple field theoretical model which plays the role of full QCD. The tractability of the model is greatly enhanced by employing non-relativistic kinematics. This choice has no bearing on the qualitative conclusions presented here. The model Hamiltonian is

$$\hat{H} = - \int d^3x \hat{\psi}_f^\dagger \tau_3 \left(m_f - \frac{\nabla^2}{2m_f} \right) \hat{\psi}_f + g \int d^3x \hat{\psi}_f^\dagger \tau_1 \hat{\psi}_f + \frac{1}{2} \int d^3x d^3y \hat{\psi}_{f'}^\dagger(\mathbf{x}) \hat{\psi}_{f'}(\mathbf{x}) V(\mathbf{x} - \mathbf{y}) \hat{\psi}_f^\dagger(\mathbf{y}) \hat{\psi}_f(\mathbf{y}). \quad (1)$$

where the field is defined in terms of particle operators as

$$\hat{\psi}_f(\mathbf{x}) = \int \frac{d^3k}{(2\pi)^3} e^{i\mathbf{k}\cdot\mathbf{x}} \begin{pmatrix} b_f^\dagger(\mathbf{k}) \\ d_f(-\mathbf{k}) \end{pmatrix} \quad (2)$$

and the τ_i are Pauli matrices acting on the quark-antiquark doublet. The flavour index f takes on two values representing a heavy quark species of mass M and a light quark species of mass m (denoted Q and q in the following).

The theory is truncated by projecting onto the lowest two Fock space sectors via the Ansatz

$$|\Psi\rangle = \int d^3r_1 d^3r_2 \phi_{QQ}(r_1 - r_2) \varphi\left(\frac{r_1 + r_2}{2}\right) b_Q^\dagger(r_1) d_Q^\dagger(r_2) |0\rangle + \int d^3r_1 d^3r_2 d^3r_3 d^3r_4 \varphi\left(\frac{u}{2}(r_1 + r_4) + \frac{v}{2}(r_2 + r_3)\right) \psi(u(r_1 - r_4) + v(r_3 - r_2)) \phi_{Qq}(r_1 - r_3) \phi_{qQ}(r_2 - r_4) b_Q^\dagger(r_1) d_q^\dagger(r_3) b_q^\dagger(r_2) d_Q^\dagger(r_4) |0\rangle. \quad (3)$$

Here the wavefunctions ϕ represent mesons, φ are center-of-mass wavefunctions which will be irrelevant in the following, and ψ is the relative two-meson wavefunction. The parameters u and v are quark mass ratios defined as $u = M/(m + M)$ and $v = m/(m + M)$. The Ansatz has been specialised to the case representing a QQ meson and the $Qq \cdot qQ$ continuum.

We shall restrict attention to the light quark loop modifications of heavy mesons (loop effects on the heavy-light mesons and light-light mesons do not change the qualitative conclusions presented here). Computing $\langle \Psi | \hat{H} | \Psi \rangle$ under these restrictions and varying with respect to ϕ and ψ yields the coupled equations

$$H_0 \phi_{QQ}(r) + \Omega(r) \psi(ur) = E \phi_{QQ}(r) \\ H_1 \psi(\rho) + u^{-3} \Omega\left(\frac{\rho}{u}\right) \phi_{QQ}\left(\frac{\rho}{u}\right) = E \psi(\rho), \quad (4)$$

where r is the distance between the quarks in the meson and ρ is the intermeson coordinate; these are related by $\rho = ur$. Additionally,

$$\Omega(\mathbf{r}) = g \int d^3x \phi_{Qq}(\mathbf{r}/2 - \mathbf{x}) \phi_{Qq}(\mathbf{r}/2 + \mathbf{x}), \quad (5)$$

$$H_0 = 2M - \frac{\nabla^2}{M} + V(r), \quad (6)$$

and

$$H_1 \psi(\rho) = \left[2M_{qQ} - \frac{\nabla^2}{m + M} + V_D(\rho) \right] \psi(\rho) + \int d^3R V_E(R, \rho) \psi(R) - E \int d^3R N(R, \rho) \psi(R). \quad (7)$$

The dependence on u emerges because quark pair creation at a point forces $r_2 = r_3$ and $r_1 - r_4 = r$ (see the argument of ψ in Eq. 3). The factor u^{-3} is due to the integral $\int d^3r \delta(\rho - ur)$ which occurs upon taking the functional derivative of $\langle \Psi | \hat{H} | \Psi \rangle$ with respect to $\psi^*(\rho)$. Its presence can be confirmed by checking hermiticity of Eq. 4. Eq. 7 describes the interactions of the heavy-light mesons and is in the form of a resonating group equation[3]. The kernels V_E and N are nonlocal exchange and normalisation kernels whereas V_D is the direct interaction obtained by convoluting the quark potential with mesonic wavefunctions. The equations are simplified further by neglecting final state interactions in the meson-meson channel. Self energy diagrams are assumed to be entirely renormalised away and a single continuum channel is used. Again, none of these simplifications qualitatively change the conclusions. The model specification is complete upon assuming an SHO interaction

$$V(r) = \frac{1}{2} kr^2 \quad (8)$$

with spring constant $k = 0.138$ GeV³ and hence $\omega_{QQ} = 240$ MeV. This fixes the heavy-light meson wavefunction yielding a transition operator of

$$\Omega(r) = g e^{-\beta^2 r^2/4}. \quad (9)$$

with $\beta^2 = \omega_{Qq} m M / (m + M)$.

Finally quark masses shall be taken to be $M = 4.8$ GeV and $m = 0.3$ GeV. The SHO strength and the quark masses were chosen so that the model is a facsimile of the upsilon system. As is typical in constituent quark models, a flavour-dependent constant is added to the Hamiltonian. In this case $\delta H_{Qq} = -0.75$ GeV, yielding a ‘B’ meson mass of $M_{qQ} = 5.4$ GeV.

In the following the $q\bar{q}$ interaction shall be called the ‘bare’ interaction. It is denoted by the solid line in Fig. 2. The screened interaction is obtained from Eq. 4 by diagonalisation in the static quark limit. The results are the dashed lines in Fig. 2. Of course the screened potentials have the expected adiabatic surface crossing at $r \approx 4.2$ GeV $^{-1}$. It is also significant that the screened potentials exhibit a modified short range behaviour. According to Lore the adiabatic screened potential is only modified near the continuum threshold, but, as shown here, the r -dependence of the transition amplitude can affect the screened potential at short range. This can have important phenomenological effects on hadron properties and can, for example, substantially modify the assumed short range behaviour of the constituent quark model.

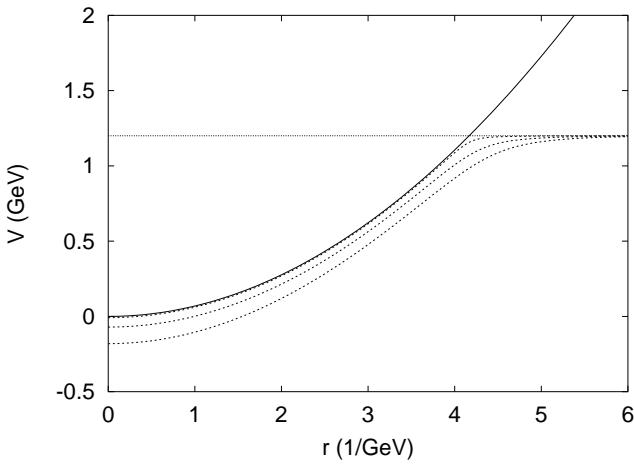


FIG. 2: The Model Screened Potential. The solid line is the bare potential; the dashed lines are the screened potential for $g = 100, 300$, and 500 MeV from top to bottom.

B. Method

It is convenient to solve the coupled channel problem by integrating out the bound state channel. The effective potential which describes the bound state physics is

$$\langle k | V_{eff} | k' \rangle = 2\pi^2 \sum_i \frac{\omega_i^*(k) \omega_i(k')}{(E - E_i)} \quad (10)$$

where E_i are the eigenenergies of H_0 and $\omega_i(k) = \langle \phi_{QQ}^{(i)} | \hat{\Omega} | k \rangle$.

The Bethe-Heitler equation is solved with the addition of a weakly coupled probe channel to locate the poles in the T-matrix. A typical computation is shown in Fig. 3. Heavy meson poles below threshold are quite narrow and easily seen by the probe channel. Identifying resonance locations above threshold is more difficult and hence Argand diagrams were often employed. This was essential in the case of the unquenched light meson spectrum because of the strong scattering which takes place in the

$(q\bar{q})(q\bar{q})$ channel. In the tables below, the real part of the poles are referred to as the ‘full’ spectrum.

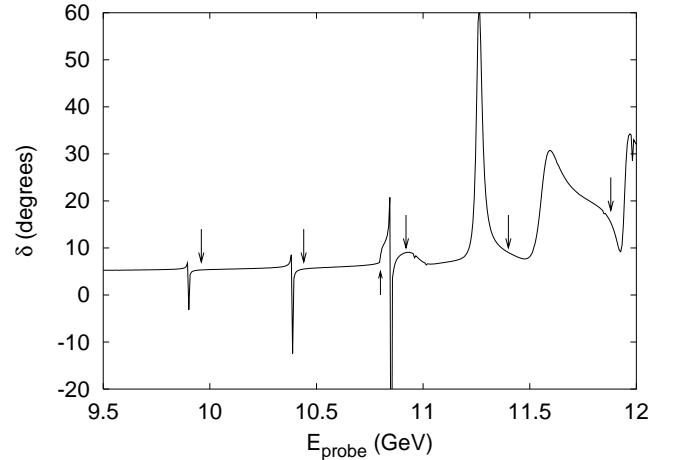


FIG. 3: Probe Phase Shift for $g = 300$ MeV. The arrows indicate bare pole locations. Threshold is shown at 10.8 GeV.

III. RESULTS AND DISCUSSION

The mass spectra for four possible solutions to the full problem of Eq. 4 are presented here:

- (i) the bare spectrum obtained by solving $H_0 \phi_{QQ} = E \phi_{QQ}$;
- (ii) the screened spectrum obtained by employing the screened potential, V_{sc} in the bare Hamiltonian;
- (iii) the renormalised spectrum;
- (iv) the full spectrum obtained by solving the full Hamiltonian, Eq. 4.

The renormalised Hamiltonian is obtained by fitting the oscillator frequency ω_R and quark mass M_R to the full (‘experimental’) low lying spectrum. This mimics the procedure followed with the quark model and serves as a point of comparison with the full model and the two alternate truncations.

As mentioned above, the use of the screened potential is frequently adopted in hadronic physics in an attempt to model the effects of coupled channels¹. The efficacy of this idea will be tested by comparing with the full spectrum below.

¹ Unfortunately, a recent attempt to use screened potentials to describe heavy mesons[4] is not a good test of the method since the ionization threshold was chosen more than 1 GeV above the physical threshold and the string tension was adjusted so that the screened potential approximated the bare potential in the physically relevant region.

A. Heavy Mesons

Model parameters were chosen to mimic the upsilon system: $M = 4.8$ GeV, $m = 0.3$ GeV and $\omega_{QQ} = 0.24$ GeV. The lightest heavy-heavy meson has a mass of 9.96 GeV, the lightest heavy-light meson is 5.4 GeV, and the ‘upsilon’ decay threshold lies between the second and third heavy-heavy states at 10.8 GeV.

TABLE I: Heavy Meson Spectra (GeV) [$g=100$ MeV]

model	E_0	E_1	E_2	E_3	E_4
bare	9.960	10.440	10.920	11.400	11.880
screened	9.952	10.428	—	—	—
renorm	9.956	10.430	10.904	11.380	11.852
full	9.956	10.430	10.91	11.4	11.887

Tables I-III present the four spectra with progressively stronger decay amplitudes. The case with $g = 100$ MeV induces small shifts in the bare spectrum, as indicated by the small renormalisation: the bare parameters are $\omega = 240$ MeV and $M = 4.8$ GeV, while the renormalised ones are $\omega_R = 237$ MeV and $M_R = 4.8$ GeV. The effects on the bare spectrum and on the renormalisation become progressively larger as the coupling increases. For $g = 300$ MeV the renormalised parameters are $\omega_R = 243$ MeV and $M_R = 4.767$ GeV. For $g = 500$ MeV one obtains $\omega_R = 252$ MeV and $M_R = 4.713$ GeV.

TABLE II: Heavy Meson Spectra (GeV) [$g=300$ MeV]

model	E_0	E_1	E_2	E_3	E_4
bare	9.960	10.440	10.920	11.400	11.880
screened	9.894	10.372	10.778	—	—
renorm	9.898	10.384	10.870	11.356	11.842
full	9.898	10.384	10.847	11.255	11.578

TABLE III: Heavy Meson Spectra (GeV) [$g=500$ MeV]

model	E_0	E_1	E_2	E_3	E_4
bare	9.960	10.440	10.920	11.400	11.880
screened	9.794	10.279	10.708	—	—
renorm	9.803	10.307	10.811	11.315	11.819
full	9.803	10.307	10.758	11.29	11.63

These results indicate that the screened potential is quite accurate below the ionisation threshold for heavy quarks. However, it is useless, and in fact, unphysical above this point. Alternatively the renormalised spectrum remains useful, even above threshold. This is true even for $g = 500$ MeV where mass shifts as large as 160 MeV are obtained. For example, upon renormalisation, the error in the binding energy of the fifth resonance is 9%. Deviations between the full model (‘experiment’) and the renormalised model (the ‘quark model’) appear once one has exhausted the available parameters and appear to get worse as one goes higher in the spectrum.

B. Light Mesons

One expects the heavy spectrum to be relatively stable under the perturbations caused by quark pair creation and this is borne out by the computation above. Alternatively, light mesons are likely to suffer greater effects. This is investigated here using the same model as above in the light-light sector. Thus $\omega_{qq} = 960$ MeV and the quark masses are $M = m = 300$ MeV. This yields a ground state meson of mass 2.04 GeV and a decay threshold of 4.08 GeV. Table IV shows the resulting spectra for $g = 100$ MeV. In this case the renormalised parameters are $\omega_R = 965$ MeV and $M_R = 286$ MeV. Increasing the coupling to $g = 300$ MeV yields the results of Table V with parameters $\omega_R = 958$ MeV and $M_R = 232$ MeV.

TABLE IV: Light Meson Spectra (GeV) [$g=100$ MeV]

model	E_0	E_1	E_2	E_3	E_4
bare	2.04	3.96	5.88	7.80	9.72
screened	2.036	3.844	—	—	—
renorm	2.02	3.95	5.88	7.81	9.74
full	2.02	3.95	5.87	7.80	9.72

In these cases the screened adiabatic potentials do not perform as well as in the heavy quark case. This can perhaps be expected since the screened potential is based on the adiabatic heavy quark limit of hadronic interactions. However, even under the large mass shifts of the $g = 300$ case, the renormalised model does quite well, with a 1.5% error in the binding energy of the fifth resonance.

TABLE V: Light Meson Spectra (GeV) [$g=300$ MeV]

model	E_0	E_1	E_2	E_3	E_4
bare	2.04	3.96	5.88	7.80	9.72
screened	2.022	3.834	—	—	—
renorm	1.901	3.816	5.731	7.646	9.561
full	1.901	3.816	5.618	7.774	9.70

C. Applications to Lattice Computations

The lessons of the previous section can be profitably applied to lattice gauge computations of adiabatic potentials and their mixing. The canonical application is to the Wilson loop potential. In particular, extensive effort has been devoted to finding ‘string breaking’ in the Wilson loop potential in unquenched QCD[5]; yet the expected surface crossing to the meson-meson continuum has been surprisingly difficult to observe. The common explanation is that better operators are required to see string breaking. But there is a simpler explanation: the confinement potential and the meson-meson continuum are connected via a Fock sector transition operator which happens to be small. Indeed, the mixed adiabatic surfaces are estimated to be split by only 50 MeV[6]. Thus surface mixing cannot be observed on the lattice unless

large temporal separations are achieved[7]. A recent lattice computation which finally observes string breaking confirms these assertions: large temporal extensions in Wilson loops were required and the calculated adiabatic potential gap is roughly 50 ± 4 MeV[8].

A second example from the lattice illustrates the importance and the subtlety of surface mixing in modelling. We consider the static energy of the first excited state of the flux tube in the quenched approximation, denoted as the Π_u adiabatic potential. This potential describes the quark dynamics of a heavy hybrid meson state. As seen in Fig. 4, the Π_u potential increases at short distance, in keeping with the expected perturbative colour octet interaction of $V = +\alpha_s/(6r)$. However, it becomes energetically favourable for the system to emit a gluon once the interquark separation becomes sufficiently small. This gluon combines with the ‘valence’ gluon of the hybrid to form a scalar glueball, thereby changing the quark colour configuration to that of a singlet, and hence the adiabatic potential to the attractive $V = -4\alpha_s/(3r)$ form. In short, it is energetically favourable for the hybrid to convert to a ground state meson and a glueball at short distances. The point at which this should happen is indicated by the arrow in Fig. 4. It is thus somewhat perplexing that the lattice data passes through this point with no transition seen.

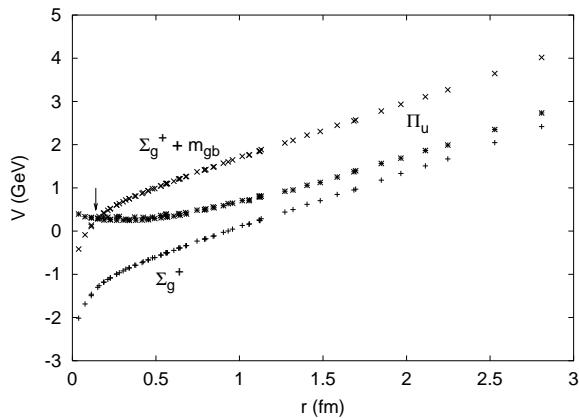


FIG. 4: Adiabatic Surface Crossing for Heavy Hybrid Potentials. The standard Wilson loop potential is denoted Σ_g^+ while the first excited potential is denoted Π_u . Lattice data are from Ref.[9].

It is possible that the expected surface mixing does not occur because the minimal relative momentum permitted on the lattices employed in the study were too coarse to permit the required hybrid decay. However, a physical reason for the suppression of this coupling is also possible. For example, if one considers the hybrid to be dominated by a Fock space component consisting of a constituent quark, antiquark, and gluon, then the postulated transition occurs through gluon emission from either the valence gluon or a valence quark. In the former case the coupling to a glueball is zero due to the colour overlap while in the latter case the coupling is suppressed

by the large (infinite) quark mass. Thus one expects the coupling between the surfaces to be very small, which implies that the surface mixing will not be seen unless lattices with exceptionally large temporal extents are employed.

While this is a plausible explanation for the observed short range behaviour of the Π_u potential it begs the question: which surface should one use when modelling heavy hybrids? The observations of the previous section dictate that one should use the unmixed lattice potential – mixing is a non-potential effect and should not be considered in valence hybrid models (in the same way that screened potentials are not useful for modelling excited canonical mesons and baryons)[10].

IV. CONCLUSIONS

The lessons gleaned from the simple mixing model presented here are directly applicable to the constituent quark model. One sees, for example, that mixing with the continuum can induce an effective short range interaction. This interaction is in general spin-dependent and can have a substantial effect on phenomenology. Sorting out spin-dependence due to gluon exchange, instantons, relativistic dynamics, and continuum mixing remains a serious challenge for model builders[11]. Another lesson concerns the lore that passing the lowest continuum threshold in a given channel is associated with a deterioration of the quality of potential models of hadrons. But we have seen that hadron masses are shifted throughout the spectrum (including below threshold) due to a given channel. Rather, it is the proximity of a continuum channel which can cause local distortions of the spectrum². Of course, the problem in hadronic physics is that continuum channels tend to get dense above threshold.

A number of additional issues complicate the application of these ideas to hadronic physics. For example, the QCD transition operator is determined by nonperturbative gluodynamics and is certainly not as simple as that of Eq. 1. In addition, as stated above, many continuum channels are present. Summing over these is nontrivial – indeed the sum may not converge (see for example, Ref. [6]). Furthermore, one expects that when the continuum virtuality is much greater than Λ_{QCD} quark-hadron duality will be applicable and the sum over hadronic channels should evolve into perturbative quark loop corrections to the Wilson loop/quark model potential. Correctly incorporating this into constituent quark models requires marrying QCD renormalisation with effective models and is not a simple task. Finally, pion and multipion loops can be expected to dominate the virtual continuum component of hadronic states (where allowed) due to the light

² Mass shifts due to a single continuum are small for bare states far from threshold.

pion mass. This raises the issue of correctly incorporating chiral dynamics into unquenched quark models. The relationship of chiral symmetry breaking to the constituent quark model has been discussed in Ref. [12] and a variety of hadronic models which incorporate chiral symmetry breaking exist[13] but much remains to be achieved.

The results presented here imply that, in analogy with the renormalised model, constituent quark models must become progressively less accurate high in the spectrum. However, this effect is likely to be overwhelmed by more serious problems: the nonrelativistic constituent quark model must eventually fail as gluonic degrees of freedom are activated and because chiral symmetry breaking (which is the pedestal upon which the constituent quark model rests) becomes irrelevant for highly excited

states[14]. It is clear that an exploration of the excited hadron spectrum is required to understand the interesting physics behind the unquenched quark model, gluonic degrees of freedom, and chiral symmetry breaking.

Acknowledgments

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